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We present a broad program for using high energy neutrino astronomy with large neutrino detectors to directly test for the existence of heavier weak intermediate vector bosons (IVB's) and cosmic antimatter. Changes in the total cross section for $\nu N \rightarrow \mu X$ due to additional propagators are discussed and higher mass resonances in the annihilation channel $\bar{\nu}_e e^- \rightarrow X$ are analyzed. The annihilation channel is instrumental in the search for antimatter, particularly if heavier IVB's exist.

I. INTRODUCTION:

The Glashow-Weinberg-Salam (GWS) model of electroweak interactions has proven to be an excellent theory for the description of all present accelerator data. However, questions have been raised as to whether GWS provides a completely fundamental description, and extended electroweak theories and composite models of quarks and leptons have recently been proposed. In general, these theories give different values for the "standard" W , Z boson masses and predict the existence of additional weak intermediate vector bosons. Therefore, the search for weak intermediate vector bosons (IVB) will provide an important test to choose between these theories. In this paper, we will connect the IVB question with the question of the existence of antimatter on an astronomical scale elsewhere in the universe. This is another fundamental question which, although controversial, must also be decided by experiment and observation.

Tests for these two important issues would not seem at first to be directly related. However, we will present here a cosmic ray neutrino astronomy program which addresses both questions¹. The key role in this program would be played by a cosmic ray neutrino telescope, such as the proposed DUMAND (Deep Underwater Muon And Neutrino Detector) facility² which has been proposed by a large collaboration of physicists and astrophysicists.

Experimental efforts to renew the search for the weak intermediate vector boson are now underway. IVB's within the mass range predicted by the GWS model are well within reach of the CERN $p\bar{p}$ collider; thus a clear answer as to their existence can be expected in the near future. However, other models (See Table I in Section II.) predict larger masses which may be beyond the reach of the $p\bar{p}$ collider. (In this regard, it should be noted that some recent data³ could be interpreted to mean that the IVB mass is greater than

100 GeV/c². This remains to be confirmed.) Thus, the search for heavier IVB's may be beyond reach of existing accelerator experiments. We should also note the great difficulty in detecting IVB's buried in the background of hadronic debris, especially if the leptonic branching ratios are small.

We suggest, as the first part of our program, that heavier bosons be sought in cosmic ray neutrino detectors through propagator effects in the charged current cross sections. We will discuss this point in detail in Section II. The energy range that can be investigated in this way goes well beyond that of present and planned accelerators. The appearance of these bosons as generalized "Glashow resonances"⁴ in $\bar{\nu}_e e^-$ annihilation can then be studied in conjunction with the propagator effects and the existence of such resonances can be used to distinguish cosmic ν_e 's and $\bar{\nu}_e$'s. The cosmic ray secondary $\bar{\nu}_e$'s will be produced in antimatter regions of the universe, as we will show in Section III. Their detection through the resonance channel thus provides the promised test for cosmic antimatter.

Since the advent of the Dirac theory of the electron and the discovery of the positron, physicists and astrophysicists have speculated as to the existence of significant quantities of antimatter on a cosmic scale in the universe. The theoretical approach to this question has changed dramatically from the original model of Alfven and Klein⁵ to the model of Omnes⁶ and recently to a qualitatively different approach based on grand unification and spontaneous symmetry breaking⁷. The basic physics argument regarding the question of a baryon symmetric versus an asymmetric cosmology hinges on the manner in which CP violation is incorporated into unified gauge theories (and into nature). If the CP violation is spontaneous, it will arise with random sign changes in causally independent regions⁷ and the universe will naturally split into domains of baryon and antibaryon excesses with no preferred

direction of CP symmetry nonconservation. Spontaneous breaking of CP leads to a domain structure in the universe with the domains evolving into separate regions of matter excess and antimatter excess. The creation of these excesses subsequent to a period of exponential horizon growth (a dynamical effect of the Higgs fields) can result in a universe in which matter galaxies and antimatter galaxies are formed in separate regions of the universe. Also, with the advent of grand unification theories, models have been suggested to generate a universal baryon asymmetry, with the consequence that no important amounts of antimatter would be left in the universe at the present time⁸. These models have been motivated by observational constraints on antimatter at least in our region of the universe⁹. However, some of these constraints have been shown to overrestrictive^{10,11} and the exciting possibility of cosmic antimatter has gained added interest from more recent observational results. Data on the cosmic background γ -radiation¹² have the spectral characteristics of cosmological $p\bar{p}$ annihilation radiation¹³. Also, recent measurements of cosmic ray antiproton fluxes^{14,15}, particularly at low energy¹⁵ where secondary antiproton production should be practically nonexistent, have suggested the possibility (among others) that we are seeing primary extragalactic antiprotons which have been accelerated to cosmic-ray energies in regions of the universe containing antimatter galaxies¹⁶. New searches for cosmic-ray antimatter are being planned¹⁷. These experiments are important. However, a completely different, independent experimental approach to this question, making use of Glashow resonances to use cosmic-ray neutrino detectors such as DUMAND to search for cosmic antimatter regions, can play a crucial role in providing an answer. This method has been suggested in connection with the GWS standard IVB¹⁸. We will see in Section III that the existence of higher mass Glashow resonances could greatly facilitate the use

of this technique. Results of a significant neutrino search for cosmic antimatter would have profound implications for our understanding of the large scale structure and evolution of the universe^{7,19}.

II. INELASTIC NEUTRINO SCATTERING

The standard GWS model is mediated by three vector bosons, W^\pm and Z^0 , which have not yet been found. If radiative corrections are included, GWS predicts²⁰

$$M_W = \frac{38.5}{\sin\theta_W} \approx 80 \text{ GeV}/c^2, \quad (2.1)$$

and (with only Higgs isodoublets)

$$M_Z = \frac{M_W}{\cos\theta_W} \approx 90 \text{ GeV}/c^2, \quad (2.2)$$

which are in agreement with the presently best lower limits²¹,

$$M_Z, M_W \gtrsim 30 \text{ GeV}/c^2. \quad (2.3)$$

These lower limits arise from the absence of propagator effects; such effects are the subject of our discussion at this stage in the program.

The effect of a finite-mass M_W on the inelastic charged-current scattering of a neutrino off of a nucleon,

$$\nu_\ell N \rightarrow \ell X, \quad (2.4)$$

is to introduce a factor $(1 + Q^2/M_W^2)^{-2}$ into the differential cross section, where $-Q^2 < 0$ is the square of the four-momentum-transfer. This changes a

linear energy dependence for the total cross section into a logarithmic growth at neutrino energies well above, producing the plateau shown in Fig. 1. This curve has been calculated by a numerical integration of the quark-parton formula

$$\frac{d^2\sigma}{dx dy} = \frac{G_F^2 ME}{\pi} \frac{1}{(1+z)^2} [q(x) + \bar{q}(x) (1-y)^2] \quad (2.5)$$

with

$$y = (E - E_\mu)/E,$$

$$x = Q^2/sy, \quad (2.6)$$

$$z = Q^2/M_W^2,$$

in terms of the Fermi constant G_F , neutrino energy E , muon energy E_μ , nucleon mass M , and c.m. energy squared $s \approx 2ME$. The quark distributions for an "average" nucleon ($1/2$ proton + $1/2$ neutron) are²²

$$\bar{q}(x) = 0.3(1-x)^7 \quad (2.7)$$

$$q(x) = \bar{q}(x) + [1.79(1-x)^3(1+2.3x) + 1.07(1-x)^{3.1}]x^{1/2},$$

It is difficult to detect a standard- W propagator effect with existing accelerators and even with future fixed-target-neutrino-beam machines. A very recent report³, however, claims enough sensitivity to place a lower limit $M_W \gtrsim 100 \text{ GeV}/c^2$ at the 90 percent confidence level. However, scaling violations and new physics are alternative explanations for this result, so

that careful confirmation needs to be made before conclusions are definitely drawn. It is hoped that the issue of standard W's and Z's will be settled soon in the proton-antiproton SPS collider, through actual production and (difficult) detection.

If additional, heavier weak bosons exist, they may be even harder to find in collider experiments. Leptonic branching ratios are expected to be small as in the standard case, and hadronic decays are buried in background debris. Also, other kinds of new particles may masquerade as W's, and the c.m. energy required for production may be beyond the range of the collider. The mass ranges for IVB's given by different electroweak alternatives can be quite large. We have listed some expectations for multi-W models in Table I.

Among these multi-W models are those^{23,24} which extend the SU(2) X U(1) GWS gauge group to SU(2) X U(1) X G where G is an arbitrary group and which are still spontaneously broken renormalizable gauge theories. The spectrum of weak bosons depends on G. Previous interest has been in the possibility that some IVB's in this theory might be lighter than the standard weak bosons [Eqs. (2.1) and (2.2)] and thus easier to detect. The possibility of heavier IVB's is of more significance here.

A more general situation regarding electroweak theories has been studied²⁵, in which agreement with GWS is required in the low energy limit, but the requirement of renormalizability is relaxed. In this case, an arbitrary number of weak bosons can be considered with global SU(2) and with mixing between a "primeval" photon and the neutral weak bosons. The mean of several charged boson masses is then constrained to be $\langle M_W \rangle < 37.4 \text{ GeV} / \sin^2 \theta_W \approx 163 \text{ GeV}$. This may arise if the W's, quarks and leptons are composite particles made up of constituents bound by "hypergluons."²⁶ This theory may be ultimately renormalizable, reflecting a gauge theory where weak

interactions become strong at high energies, a strong-coupling confining version²⁷ of GWS. In such a gauge theory or quantum hypercolor dynamics (QHD), with confinement parameter $\Lambda_H \sim G_F^{-1/2}$, the weak interaction is an effective "Van der Waals" interaction generated by QHD and is analogous to the description of the strong nuclear force by QCD quark-gluon theory. The relation $g^2/8M_W^2 = G_F/\sqrt{2}$ would then imply $M_W = 123g \text{ GeV}$ with $g \sim 1$. The possibility of weak bosons which couple to right-handed currents has also been much discussed²⁸.

The emphasis in these multi-W deliberations has been on the role of the lighter bosons in experiments. To test for the existence of heavier partners, particularly in the context of our overall program, and noting the aforementioned difficulties, we may turn to cosmic-ray neutrinos and thus return to the propagator effect search in reaction (2.4). The propagator effect can be restated in terms of the y distributions shown in Fig. 2. It has been argued that an experiment using cosmic-ray neutrinos would be successful in detecting the small y enhancement²⁹. Neither an antineutrino admixture as large as 25 percent nor the asymptotic freedom corrections seems to mask the y effect.

We may ignore real production of W's in neutrino interactions at these higher energies. The reason is as follows. The production cross sections for

$$\nu_e + N \rightarrow e^- + W^+ + X \quad (2.8)$$

have been recalculated³⁰. They are shown in Figure 1. Although real production dominates charged current scattering (2.4) for $M_W \lesssim 10 \text{ GeV}/c^2$ (now ruled out by experiment), the reverse is true in the unified electroweak models discussed here for which $M_W \sim 100 \text{ GeV}/c^2$ or greater. This is due to the fact that $G_F^2 M_W^2$ sets the scale for the deep inelastic cross section

shoulder and G_F sets the scale for the semi-weak production (whose threshold is at the same place as the corresponding shoulder). The larger the value of M_W nature chooses, the larger the ratio of the "charged-current" hadron production to the W production.

We now consider what changes arise from a sequence of left-handed W's, W_1 with mass M_1 . For the moment, we restrict ourselves to a pair of W's, W_1 and W_2 , which is sufficient to analyze the multi-propagator effect. We define

$$\kappa \equiv \left(\frac{M_2}{M_W}\right)^2, \quad \epsilon \equiv 1 - \left(\frac{M_1}{M_W}\right)^2 \quad (2.9)$$

and we assume that the lighter W_1 is close to the standard single W- particle mass and that W_2 has a larger mass:

$$M_1 \approx M_W, \quad M_2 \gtrsim 3M_W \quad (2.10)$$

so that

$$\kappa \gg 1, \quad \epsilon \ll 1$$

We see from the references^{23,24} on multi-W and multi-Z alternatives to the standard GWS model that it is consistent with the overall charged-current strength to take the leptonic decay widths to be

$$\frac{\Gamma_1^\ell}{M_1} \approx \frac{\Gamma_W^\ell}{M_W} \approx .003 \quad (2.11)$$

$$\frac{\Gamma_2^L}{M_2} \approx \epsilon \kappa \frac{\Gamma_W^L}{M_W} . \quad (2.12)$$

For numerical work, a value for $\epsilon \kappa$ in equation (2.12) is needed. If we assume that

$$\epsilon \kappa \approx 1, \quad (2.13)$$

then the relative couplings W_1 and W_2 are roughly comparable. Only for $\epsilon \kappa \gtrsim 10$ does the W_2 leptonic decay width become uncomfortably large. (Note that $\Gamma = 12 \Gamma_W^L$ for the usual single-boson 3-generation model.)

The standard propagator is now replaced by two terms:

$$\frac{1}{1+z} \rightarrow \frac{1}{1+z} + \frac{1}{\kappa+z} \approx \frac{1}{1+z} \frac{2z+\kappa}{z+\kappa} \quad (2.14)$$

Noting that $\kappa = \infty$ corresponds to the standard model, we have plotted the deviation of $\sigma(\kappa)/\sigma(\infty)$ from unity in Fig. 3 for $\kappa = 10$. (This deviation scales as s/M_W^2 . We see that the first plateau occurring at $s/M_W^2 \approx 1$ is pushed to a second plateau beginning around $s/M_W^2 \approx \kappa$ which is ~ 30 percent higher.

An increase in κ for fixed ϵ can enhance the deviation from the standard result. The consequence of adding more W 's to the sequence depends on the model but typically corresponds to smaller increases on up the line. (Each plateau is only a fractional increase on the previous increase.) These effects are not significant enough to be shown.

We have also calculated the y -distributions and found the same sequence of plateaus for any given value of y . Thus a plot of $([\frac{d\sigma}{dy}(\kappa)]/[\frac{d\sigma}{dy}(\infty)]) - 1$

III. NEUTRINO-ELECTRON ANNIHILATION

The search for the propagator-plateau effect can be augmented by a concomitant search for $\bar{\nu}_e e^-$ resonances³¹. The "Glashow resonance",

$$\bar{\nu}_e + e^- \rightarrow W^-, \quad (3.1)$$

was studied long ago⁴ and more recently has been the subject of a proposal for a DUMAND cosmic-ray experiment¹⁸. Atmospheric $\bar{\nu}_e$'s (from the prompt decay of heavy particles) with energies in the neighborhood of the resonance energy,

$$E_W = 6.26 \frac{s}{M_W^2} \text{ PeV} \quad (1 \text{ PeV} = 10^{15} \text{ eV}), \quad (3.2)$$

$$s = 2mE,$$

may annihilate up to 10 atomic electrons per year in the proposed detector, producing showers. The number of resonance-induced showers could be much larger if additional cosmic $\bar{\nu}_e$ fluxes exist.

The total resonance cross section for a given boson W_i with mass M_i is

$$\sigma_i \equiv \sigma(\bar{\nu}_e e^- \rightarrow W_i^- + \text{all}) = 24\pi \frac{\Gamma_i \Gamma_i^l}{(s-M_i^2)^2 + \Gamma_i^2 M_i^2} \quad (3.3)$$

where Γ_i is the total width and Γ_i^l is the leptonic channel partial width. It may be assumed that the energy dependence of the flux and detector are negligible over the widths of the resonances, ΔE_i (we assume that $\Delta E_i \ll M_i$). Therefore, the integrated cross section is

$$\int \sigma_i ds = 2m \int \sigma dE = 24\pi^2 \frac{\Gamma_i^l}{M_i} . \quad (3.4)$$

QCD corrections to the three generation result $\Gamma \approx 12 \Gamma_L$ give a more precise relation for the standard boson width²⁰

$$\Gamma \approx 9 \Gamma_L \left(1 + \frac{\alpha_s}{\pi}\right) + 3 \Gamma_L \approx 12.43 \Gamma_L . \quad (3.5)$$

with the corresponding standard boson leptonic width

$$\Gamma_L = \frac{G}{\sqrt{2}} \frac{M_W^3}{6\pi} = 0.226 \text{ GeV} \quad (3.6)$$

so that

$$\int \sigma_W ds = 2.58 \times 10^{-28} \text{ cm}^2 \text{ GeV}^2 . \quad (3.7)$$

From (3.4) we see that heavier bosons can be produced at comparable rates for a given flux and detection efficiency. If the detectors are increasingly sensitive to higher energies and if there is any plateau or shoulder in the flux (as will be discussed in the next sections on cosmic antimatter tests), the number of heavier W_i events can grow with i .

There are background showers expected from deep inelastic scattering by any neutrino flux component present at a given energy of interest. The background event rate per electron from $\sigma_{\nu N}$ integrated over the width of the resonance energy

$$\Delta E_i = \frac{M_i \Gamma_i}{2m} \quad (3.8)$$

provides the "noise" over which the resonance "signal" must be seen.

The signal-to-noise ratio is then

$$R \equiv \frac{\int_{\Delta E_i} \sigma_i dE}{\int_{\Delta E_i} \sigma_{vN} dE} = \frac{\frac{1}{2m} \int \sigma_i ds}{\langle \sigma_{vN} \rangle \frac{M_i \Gamma_i}{2m}} = \frac{24\pi^2}{\langle \sigma_{vN} \rangle} \left(\frac{\Gamma_i^2}{\Gamma_i} \right) / M_i^2. \quad (3.9)$$

For general i ,

$$\kappa_i \equiv (M_i/M_W)^2 \quad (3.10)$$

$$\eta_i \equiv 12\Gamma_i^2/\Gamma_i.$$

and using

$$\langle \sigma_{vN} \rangle \approx 3 \times 10^{-34} \text{ cm}^2, \quad (3.11)$$

$$2\pi^2/M_W^2 \approx 1.2 \times 10^{-30} \text{ cm}^2,$$

we have

$$R = 3.6 \times 10^3 \eta_i / \kappa_i. \quad (3.12)$$

Thus, the signal-to-noise ratio can be good even for large κ_i .

The ability to see a heavier W_i in deep inelastic events would allow us to assume its existence as a resonance channel and to provide a measure of the $\bar{\nu}_e$ flux. This will play an important role as a test for cosmic antimatter.

IV. LOOKING FOR AN ANTIMATTER SIGNATURE IN THE DIFFUSE COSMIC NEUTRINO BACKGROUND

In order to discuss the possibility of looking for an antimatter signature in the diffuse cosmic neutrino background, we will draw heavily on calculations of diffuse cosmic neutrino fluxes reported previously³². A production mechanism of particular importance in this context because of its large inherent charge asymmetry involves the photoproduction of charged pions by ultrahigh energy cosmic rays interacting with the universal 3K blackbody background radiation. The most significant reactions are



which occur in the astrophysical context principally through the resonance channels



because of the steepness of the ultrahigh energy cosmic ray spectrum³³.

The principal charged pion decay modes are, of course,



$$\begin{aligned} \pi^- &\rightarrow \mu^- + \bar{\nu}_\mu \\ &\quad \downarrow \\ &\quad e^- + \nu_\mu + \bar{\nu}_e \end{aligned} \quad (4.4)$$

The four leptons resulting from pion decay split the pion rest energy almost equally and there is no asymmetry in ν_μ versus $\bar{\nu}_\mu$ production. However, the π^+ decays produce ν_e 's whereas the π^- decays produce $\bar{\nu}_e$'s. Thus, if one can distinguish ν_e 's and $\bar{\nu}_e$'s in one's detector, in principle the diffuse neutrino background can tell us the ratio of ultrahigh energy protons to antiprotons in the universe. (The universe is transparent to ν_e 's and $\bar{\nu}_e$'s coming from all observable distances.) It has been pointed out that there is a significant and potentially useful way of distinguishing $\bar{\nu}_e$'s from ν_e 's, namely through their interactions with electrons^{18,31}. The $\bar{\nu}_e$'s have an enhanced cross section through the formation of weak intermediate vector bosons discussed previously in Sec. III, and we may build on the groundwork laid in Sec. III for the resonance formation of higher mass intermediate vector bosons, W_i^- with masses M_i

$$\bar{\nu}_e + e^- \rightarrow W_i^- \quad (4.5)$$

and resonance energies

$$E_i = M_i^2/2m = 6.26 \kappa_i \text{ PeV}. \quad (4.6)$$

The cosmic and atmospheric fluxes for $\bar{\nu}_e$'s, based on the calculations in Reference 32, are shown in Fig. 4.

The atmospheric $\bar{\nu}_e$ flux from π and K decay have been estimated based on

measured atmospheric muon fluxes. The dominant atmospheric contribution is expected to be from prompt decay of charmed mesons. The spectrum is steep and approximates the cosmic ray spectrum. Regarding reactions (4.1)-(4.4), we note some important considerations regarding an approximate source function for the production of neutrinos in photomeson interactions with microwave blackbody photons. Let us assume all the photons to be at the average energy $\epsilon_0 = 2.7kT \approx 6.4 \times 10^{-4}$ eV so that

$$n_{bb}(\epsilon) = n_{bb} \delta(\epsilon - \epsilon_0), \quad n_{bb} \approx 400 \text{ cm}^{-3}. \quad (4.7)$$

The energy of the photon in the cosmic ray proton rest system is

$$\epsilon' = (E_p/M_p) \epsilon_0 (1 - \cos \theta). \quad (4.8)$$

There is a large peak in the photomeson production cross section at $\epsilon' \approx 0.35 M_p$ due to the Δ resonance; and since the cosmic-ray spectrum drops off rapidly with increasing energy, most of the pion production occurs at this resonance energy. Thus, from equation (4.8) we make the approximation

$$\sigma(\epsilon') = \sigma_0 \delta \left[\chi - \frac{0.35 M_p^2}{\epsilon_0 E_p} \right] \equiv \sigma_0 \delta(\chi - E_0/E_p), \quad (4.9)$$

where

$$\chi \equiv 1 - \cos \theta, \quad E_0 \equiv \frac{0.35 M_p^2}{\epsilon_0} \approx 4.8 \times 10^{11} \text{ GeV} \quad (4.10)$$

and

$$\sigma_0 \approx 2 \times 10^{-28} \text{ cm}^2.$$

Then the source function for neutrino production from $\pi \rightarrow \mu\nu$ decay may be written in the form

$$q(E_\nu) = 4\pi/2n \int_{E_\nu/2n}^{\infty} \frac{dE_\pi}{E_\pi} n_{bb} \int_0^2 x dx \int_0^x dE_p I(E_p) \sigma(E_p, x) f(E_\pi/E_p). \quad (4.11)$$

We assume all the pions to be produced at the average energy

$$\langle E_\pi \rangle = \frac{E_p}{2} \left(\frac{m_\pi^2 + 2\epsilon' M_p}{M_p^2 + 2\epsilon' M_p} \right). \quad (4.12)$$

If we use equation (5.11), equation (5.12) further reduces to

$$\langle E_\pi \rangle = \frac{E_p^2}{2} \frac{m_\pi^2 + 0.7M_p^2}{1.7M_p^2} \approx E_p/5, \quad (4.13)$$

so that the distribution function is approximated by

$$f(E_\pi/E_p) \approx \delta(E_\pi - E_p/5). \quad (4.14)$$

If we further specify the differential cosmic-ray proton spectrum by a power-law of the form $I(E_p) = K_p E_p^{-\Gamma}$, equation (4.11) reduces to

$$\begin{aligned} q(E_\nu) &= 2\pi n_{bb} \sigma_0 K_p / n \int_{E_\nu/2n}^{\infty} dE_\pi / E_\pi \int_0^2 x dx \int_{(E_0/2)}^{\infty} dE_p E_p^{-\Gamma} \delta(x - E_0/E_p) \delta(E_\pi - E_p/5) \\ &= 10\pi \sigma_0 n_{bb} K_p E_0 / n \int_{\max[E_0/2, 5E_\nu/2n]}^{\infty} dE E^{-\Gamma-2}. \end{aligned} \quad (4.15)$$

The solution to equation (4.15) may then be written in the simple form

$$\begin{aligned} q(E_\nu) &= Q_\nu, & E_\nu < E_c, \\ &= Q_\nu (E_\nu/E_1)^{-(\Gamma + 1)}, & E_\nu > E_c, \end{aligned} \quad (4.16)$$

where $E_c \equiv 5E_0/\eta \approx 2 \times 10^{10}$ GeV and

$$Q_\nu = 2.35 \times 10^{-23} K_p (E_0/2)^{-\Gamma} / (\Gamma + 1). \quad (4.17)$$

It follows from Equation (4.16) that below 10^4 PeV, the differential neutrino spectrum from reactions (4.1)-(4.4) is roughly constant and the integral production spectrum $q(> E_\nu)$ is logarithmically decreasing with energy and may also be taken to be constant. Thus, in contrast to the atmospheric $\bar{\nu}_e$ spectrum, the cosmic $\bar{\nu}_e$ spectrum will be very flat at high energies, making separate detection feasible!

Assuming that there is no significant enhancement in the flux from production at high redshifts, the integral $\bar{\nu}_e$ spectrum from $\gamma\bar{p}$ interactions is expected to be roughly constant at 10^{-18} to 10^{-17} ν_e 's $\text{cm}^{-2} \text{sr}^{-1}$ up to an energy of $\sim 2 \times 10^7$ TeV (2×10^{19} eV) above which it is expected to drop steeply. Fig. 4 shows the estimated upper limit (UL) and lower limit (LL). It is expected that the largest competing background flux of $\bar{\nu}_e$'s will be prompt $\bar{\nu}_e$'s from the decay of atmospherically produced charmed mesons. The estimated upper and lower limits for this flux are also shown in Fig. 4 and the position of the W^- resonance is indicated by an arrow. It can be seen that a cosmic $\bar{\nu}_e$ signal may be heavily contaminated by prompt atmospheric $\bar{\nu}_e$'s at the resonance energy E_W . The cosmic flux is expected to dominate the

higher energies so that the existence of higher mass bosons W_1 may be critical to any proposed test for cosmic antimatter using diffuse fluxes. In the following discussion, we will assume that such bosons exist. (It may be possible to test for their existence independently using DUMAND (see previous section) or future colliding beam accelerators.

The event rate expected for $\bar{\nu}_e$ induced W_1^- events is quite low using the "conservative" estimates for the $\bar{\nu}_e$ flux shown in Fig. 4. For example, with

$$I_{\bar{\nu}_e}(\bar{p}_\gamma) = 10^{-26} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} \quad (4.18)$$

and with [See Sec. III]

$$\int_{-r/2}^{r/2} \sigma_1 dE = \frac{1}{2m} \int \sigma_1 ds \approx 3 \times 10^{-24} \text{ cm}^2 \text{ GeV} \quad (4.19)$$

(where we have assumed a factor of 10 increase over $\int \sigma_W dE$), we find an event rate

$$r_1 = 4\pi I_{\bar{\nu}_e} N_e \int \sigma_1 dE \approx 1 \text{ event/yr} \quad (4.20)$$

for a 10^{11} ton acoustic detector. Acoustic detectors can be much more efficient at ultrahigh energies than optical detectors³⁴. However, two points may be noted regarding this low event rate: 1) It may be possible that $I_{\bar{\nu}_e}$ is significantly higher (perhaps $\sim 10^{-25} \text{ cm}^2 \text{ s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1}$) due to cosmic ray production at high redshifts³⁵. 2) No significant signal is expected otherwise unless there is significant $\bar{\nu}_\mu \leftrightarrow \bar{\nu}_e$ mixing. The probability of such mixing between two states is

$$P(\nu_1 \rightarrow \nu_2) = \sin^2 2\alpha \sin^2 (1.27 \Delta(\text{eV}^2) \frac{L(\text{m})}{E_\nu(\text{MeV})}) \quad (4.21)$$

For typical $10^{10} M_e \bar{\nu}_\mu$'s at a cosmological distance of $10^{26} m$, mixing will occur for $\Delta > \Delta_{\min}$ where $\Delta_{\min} \sim 10^{-16} \text{eV}^2$. Although present experiments are not consistent with large mixing they do not rule it out for $\Delta \lesssim 1 \text{eV}^2$.

Present data indicate³⁶

$$P = \bar{\nu}_e / \bar{\nu}_\mu \lesssim (2 \pm 7) \times 10^{-4} \quad (4.22)$$

and $\Delta < \sim 0.9 \text{eV}^2$ for maximal mixing. If $\Delta > 0.9 \text{eV}^2$, as one would expect if $m_\nu \sim 10 \text{eV}$ from cosmological considerations, and for a $\Delta \sim m_\nu^2$ and an expected mass hierarchy for ν 's similar to that for charged leptons, then from (4.22), the mixing angle is sufficiently small so that our program will be valid.

Owing to the very low probability for helicity flipping³⁷, $\bar{\nu}_e$ from any $\nu_e \leftrightarrow \bar{\nu}_e$ oscillations, if they occur, will not produce a significant resonance signal. This is because left-handed ν_e 's have the wrong helicity for the formation of a W_1^- which couples to a V-A charged current. (An interesting related point is that such left-handed ν_e 's could produce the "right-handed" W_R 's of Table 1 which mediate V+A charged currents. We could thus also rephrase our program in terms of W_R sequences.

An acoustic deep underwater neutrino detector may provide the best hope for testing for cosmic antimatter by studying the diffuse background neutrinos. The practical threshold for such devices³⁴ appears to be in the neighborhood of $10^3 - 10^4 \text{TeV}$, where the W^- resonance occurs. For higher mass resonances W_1^- , the relevant neutrino resonance energy $E_1 \approx M_1^2$ and the effective detection volume $V_{\text{eff}} \approx M_1^6$. Considering that the incident flux is expected to be roughly constant up to energies $\sim 2 \times 10^7 \text{TeV}$, one gains much in looking for higher mass Glashow resonances at higher energies. Acoustic

detectors of effective volume $\gg 10 \text{ km}^3$ (10^{10} tons) may be economically feasible and consequently event rates of $\sim 10^2 - 10^4 \text{ yr}^{-1}$ may be attained in time.

V. LOOKING FOR ANTIMATTER SIGNATURES IN COSMIC POINT SOURCES

The asymmetry in the production of charged pions in matter versus antimatter sources is reflected in cosmic-ray pp and $\bar{p}\bar{p}$ interactions as well as $p\gamma$ and $\bar{p}\gamma$ interactions. Through the principal decay modes [Eqs. (4.3)-(4.4)], this asymmetry is again reflected in a $\nu_e - \bar{\nu}_e$ asymmetry, and thus in the characteristics of events produced in deep underwater neutrino detectors. For ν -sources, these effects may be measurable at energies ~ 1 -10 TeV with optical detectors. The details of this possibility have been discussed by Learned and Stecker³⁸.

The possibility that $p\gamma$ and $\bar{p}\gamma$ interactions in sources would produce significant fluxes of $\bar{\nu}_e$'s, detectable through the W^- resonance, has been suggested by Berezhinsky and Ginzburg¹⁸ as a way of looking for cosmic antimatter. Hopefully, this interesting suggestion will be explored in more detail as our understanding of the nature of cosmic ray production in compact objects increases. The relevant interactions here would involve $\sim 10^5$ TeV cosmic rays and ultraviolet photons in sufficient quantities.

Another possible $\nu_e - \bar{\nu}_e$ asymmetry which may provide a future test for cosmic antimatter involves lower energy (5-30 MeV) neutrinos produced during the gravitational collapse of astrophysical objects. Neutrinos from gravitational collapse events may exhibit $\nu_e - \bar{\nu}_e$ asymmetries³⁹ which can be used to determine whether the collapsing object consists of matter or antimatter by separately determining the fluxes of ν_e 's and $\bar{\nu}_e$'s. However, the bursts expected from a stellar collapse in a neighboring supercluster will

be 10^8 times weaker than the $\bar{\nu}_e$ burst previously reported⁴⁰, making detection of extragalactic antimatter collapses very difficult unless the masses involved are on a much larger scale.

VI. CONCLUSION

Neutrino telescopes can be used to search for high mass IVBs and to distinguish between matter and antimatter sources of cosmic neutrinos and thus provide a direct test of baryon symmetric cosmologies. Perhaps the most promising form of the cosmic antimatter test may lie in studies of ultrahigh energy photomeson-produced neutrinos using acoustic detectors and making use of $\bar{\nu}_e e^-$ resonances. A two stage program is suggested in which the existence of higher mass W_1^- resonances is first independently established, possibly using propagator effects.

ACKNOWLEDGMENTS

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Table 1. The ranges of W-boson masses for various models of electroweak interactions.

<u>Model</u>	<u>IVB Mass (GeV/c²)</u>	<u>References</u>
local SU(2) X U(1) (standard)	$M_W \approx 80$	GWS
local SU(2) X U(1) X G (extended)	$\langle M_W \rangle \lesssim 163$	22,23
global SU(2), photon mixing	$\langle M_W \rangle \lesssim 163$	24
composite	$M_W \gtrsim 100$	25,26
left-right symmetric	$M_W \approx 3 M_W$	27

FIGURE CAPTIONS

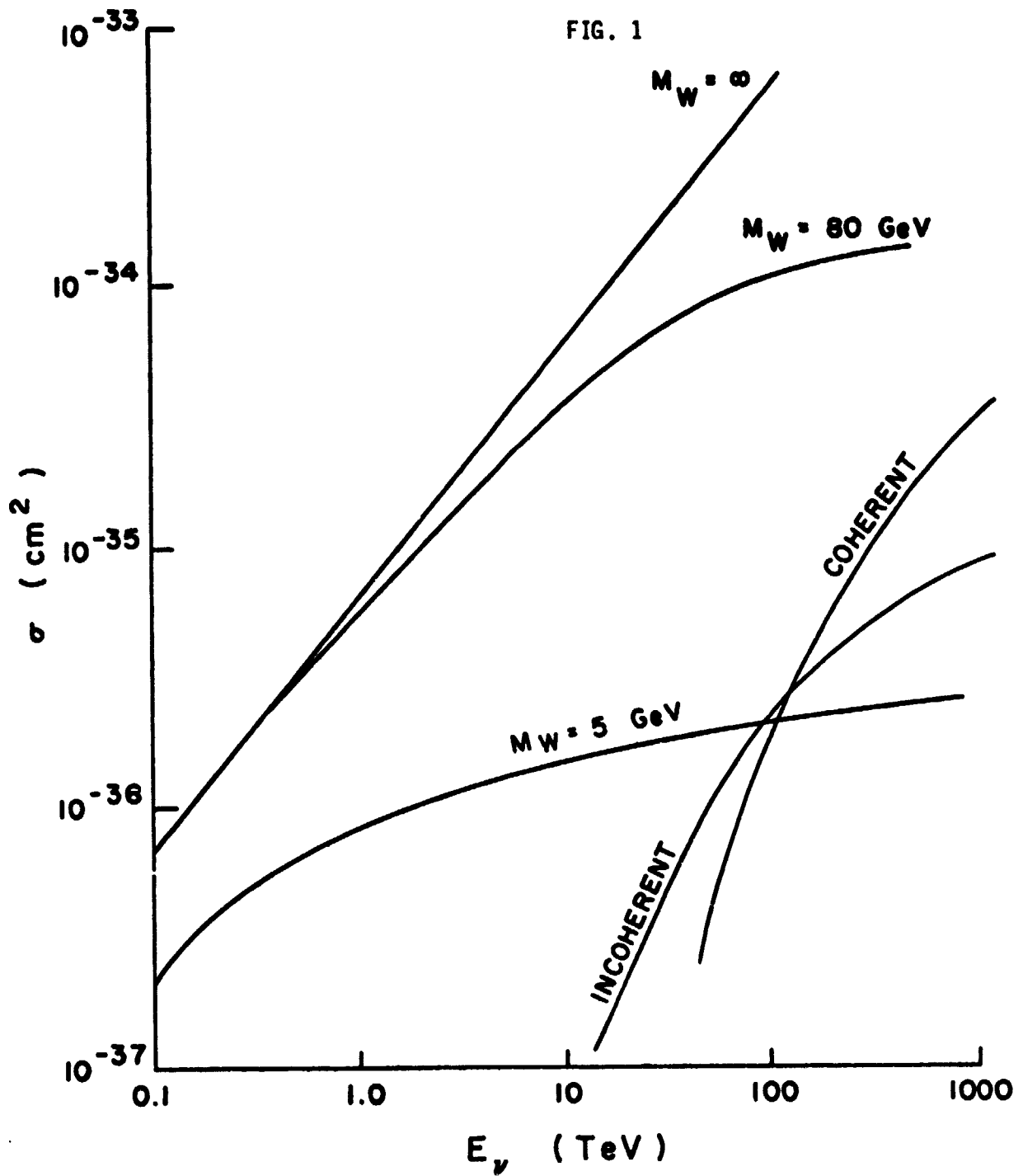
Fig. 1. Neutrino total cross sections for $\nu N \rightarrow \mu X$ where N is an average nucleon at rest. Only one weak boson is assumed with mass $M_W = 5, 80, \text{ or } \infty$. The other two curves represent real W production for $M_W = 80$. See Reference 29.

Fig. 2. The y distribution for $\nu N \rightarrow \mu X$. The solid lines show the scaling violation due to the boson propagator. See Reference 28.

Fig. 3. The increase in the total cross section for $\nu N \rightarrow \mu X$ due to a second weak boson with $\kappa = 10$. The asymptotic value is 0.36.

Fig. 4. Cosmic and atmospheric $\bar{\nu}_e$ fluxes.

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FIG. 2

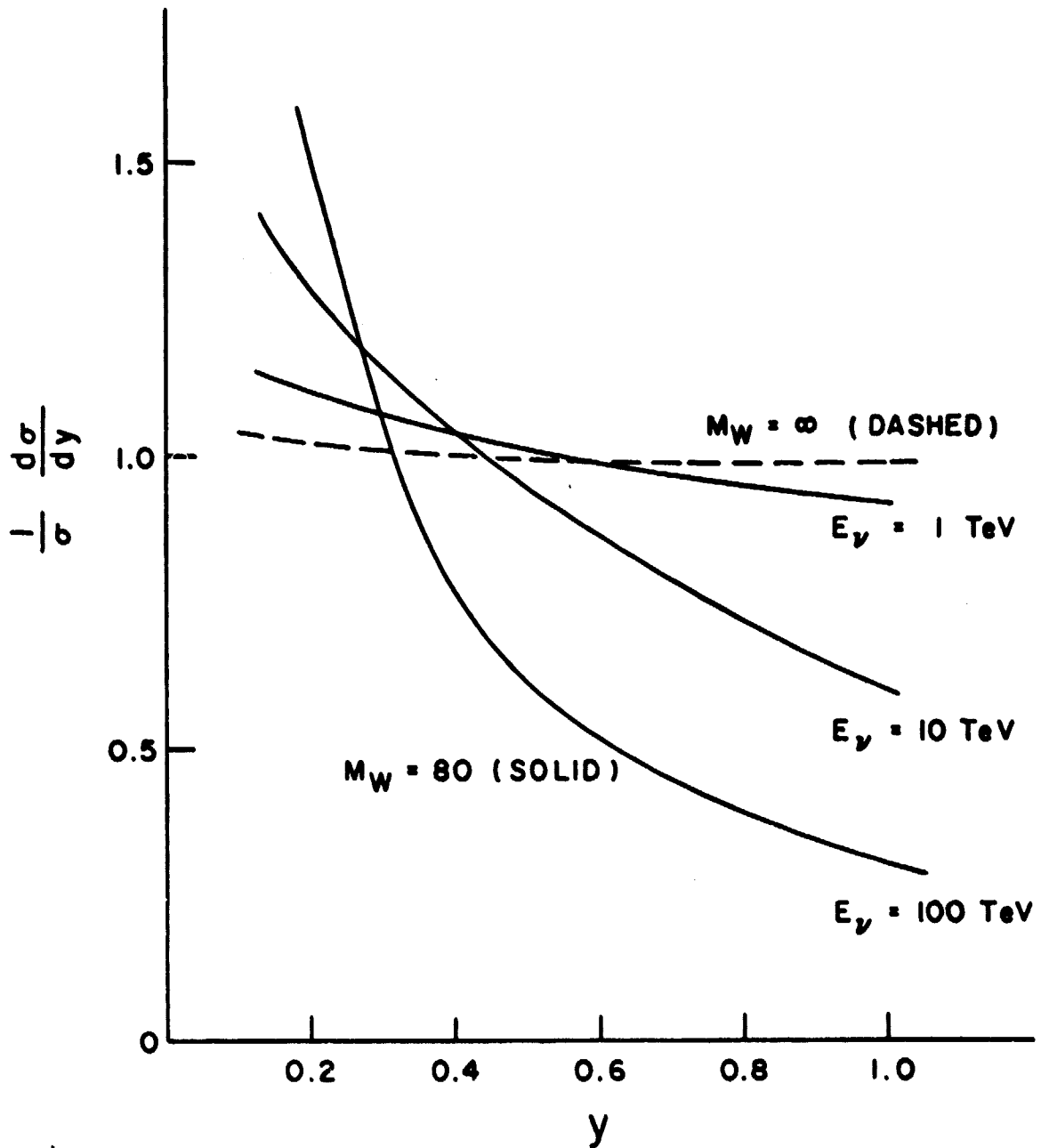


FIG. 3

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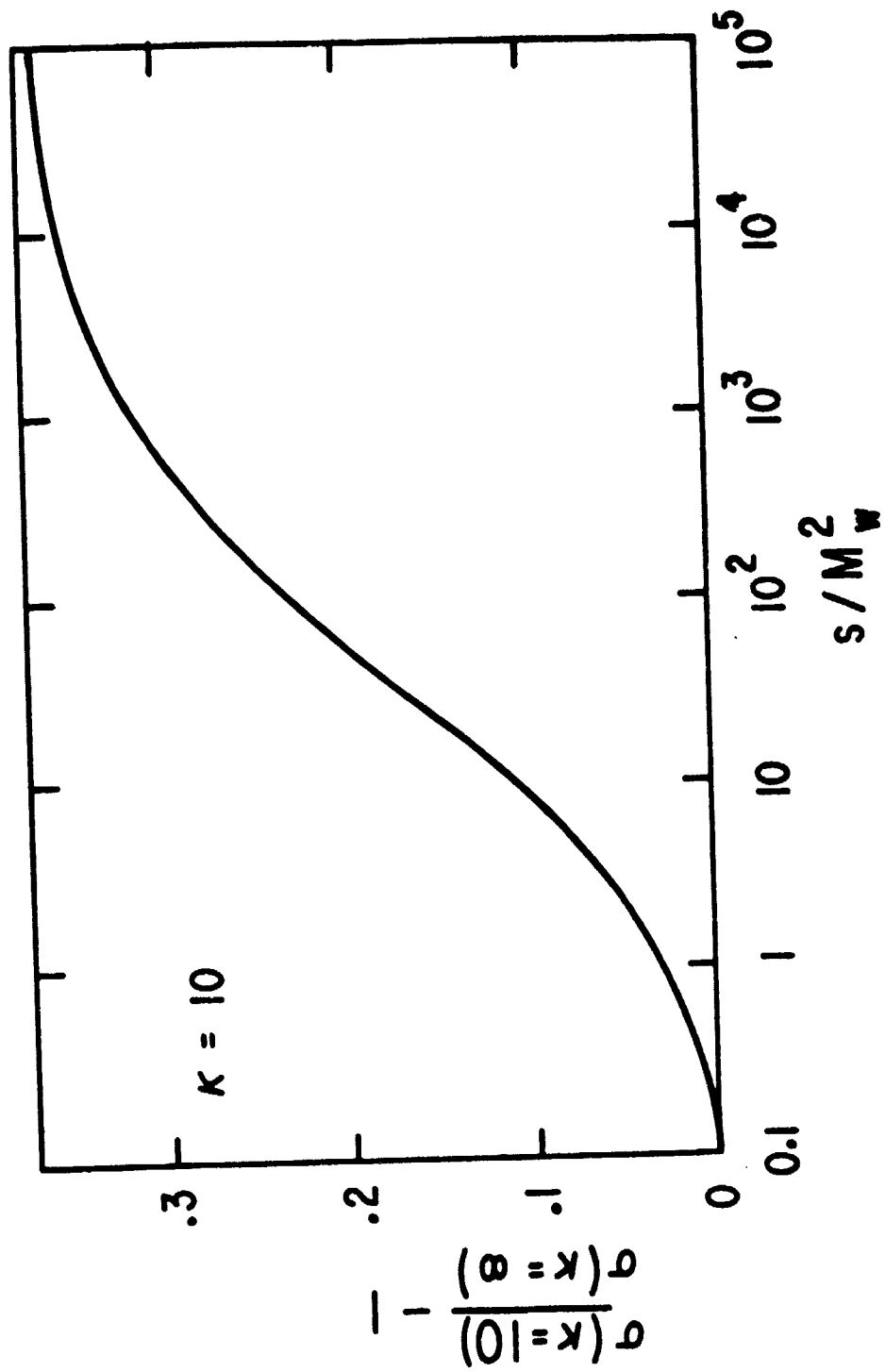


FIG. 4

